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Inelasticity effects in $\pi\pi$ production in $\gamma\gamma$ collisions*

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Abstract. The inelasticity effects in the production of pion pairs in the process $\gamma\gamma\rightarrow\pi\pi$ for real photons are investigated using the partial wave dispersion relations. The total cross sections for different photon helicities are calculated. It is observed that this process is dominated by the $\pi\pi$ final state interaction. A prediction for S* (997) $\rightarrow\gamma\gamma$ decay width is also made.

Keywords. Colliding beams; two-photon process; partial wave dispersion relations; decay width.

1. Introduction

With the advent of colliding beam experiments, there have been extensive discussions on hadronic final states in e^+e^- colliding beams via two-photon-process (Brodsky 1973, Terazawa 1973). By a two-photon process we mean a process of the type

$$e^+ + e^- \rightarrow e^+ + \gamma^* + e^- + \gamma^* \rightarrow e^+ + e^- + X \tag{1}$$

in which the colliding particles emit virtual (space-like) photons (γ^*s), which, in turn, annihilate producing a final state X, where X may be a lepton state such as e^+e^-, $\mu^+\mu^-$, or any positive charge conjugation state, such as $\pi^+\pi^-$, $\pi^+\pi^-\pi^0$, π^0, η, K^+K^-, etc. This proceeds predominantly according to the direct diagram given in figure 1(a). For the dominant part of the cross section, where both leptons scatter in the near forward directions, the equivalent photons are nearly on the mass-shell and the hadronic aspects of their interaction are expected to dominate. The measurements of the total cross section σ ($\gamma\gamma\rightarrow$hadrons) are of comparable importance to the total photoabsorption cross section and other hadronic cross section. A study of a specific channel like

$$\gamma\gamma\rightarrow\pi\pi \tag{2}$$

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allows a unique probe of the pion Compton amplitude and can determine the $\pi\pi$-scattering lengths as also the phase shifts (Carlson *et al* 1972). The behaviour of this amplitude at high energy and large momentum transfer is an important test of scaling laws and fixed-pole behaviour, predicted by Quark-Parton models (Walsh 1973, Feynman 1972).

It is the aim of the present paper to investigate the reaction (2) for real photons and for both the charged as well as neutral pion production. A partial-wave dispersion relation calculation is done to compute the cross-sections for the above process when photon helicities are (i) same and (ii) opposite. The details of calculation are presented in the subsequent sections.

2. Kinematical considerations

The process under consideration is depicted in figure 1(*b*). We define the invariants

$$\begin{aligned} s &= (k_1 + k_2)^2 = (q_1 + q_2)^2 \\ t &= (q_1 - k_1)^2 = (q_2 - k_2)^2 \\ u &= (q_2 - k_1)^2 = (q_1 - k_2)^2, \end{aligned} \tag{3}$$

with

$$s + t + u = 2. \tag{4}$$

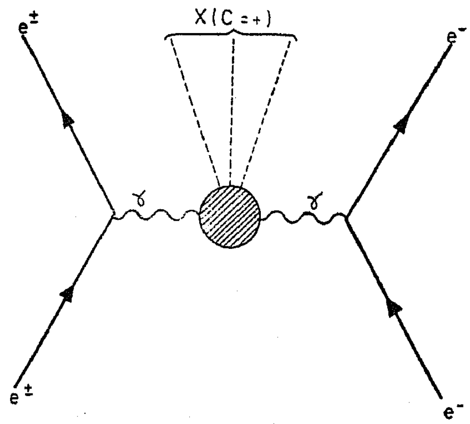

Figure 1(a). Production of a $C=+$hadron state via two photons in e^+-e^- colliding beams.

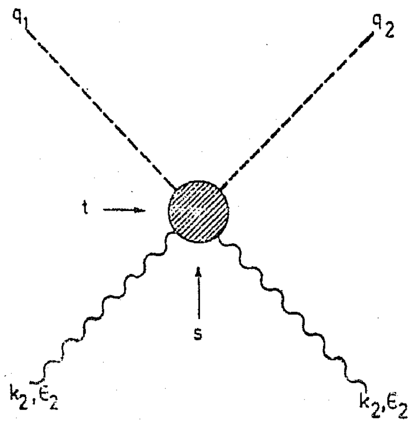

Figure 1(b). The process $\gamma\gamma\rightarrow\pi\pi$.

In the centre of mass (c.m.) system where $|\mathbf{k}_1| = |\mathbf{k}_2| = k$ and $|\mathbf{q}_1| = |\mathbf{q}_2| = q$, we have

$$\begin{aligned} q &= (s/4-1)^{1/2}, \\ s &= 4\omega^2 = 4k^2, \\ t &= 1 - s/2 + \sqrt{s}\, q\, x, \\ u &= 1 - s/2 - \sqrt{s}\, q\, x, \end{aligned} \tag{5}$$

where $x = \cos\theta$, θ being the c.m. scattering angle. Also we have the following relations amongst the invariants and the scattering angle:

$$\begin{aligned} (t-1)(u-1) &= s^2/4 - sq^2x^2, \\ (tu-1) &= sq^2(1-x^2). \end{aligned} \tag{6}$$

The amplitude for double pion production for either isoscalar photons or isovector photons can be written as (Gourdin *et al* 1960)

$$T = N^{\mu\nu}\, \varepsilon_{1\mu}\, \varepsilon_{2\nu}, \tag{7}$$

where

$$N^{\mu\nu} = \left\{ g^{\mu\nu} - \frac{k_2^\mu k_1^\nu}{k_1 \cdot k_2} \right\} T^{(1)} + \{k_2^\mu k_1^\nu (t-1)(u-1) + 2q_1^\mu q_2^\nu (u-1) + 2q_2^\mu q_1^\nu (t-1)\}\, T^{(2)}. \tag{8}$$

It has been shown (Abarbanel *et al* 1968) that the following amplitudes:

$$\begin{aligned} M_{++} &= \frac{T_{++}}{s} = \frac{1}{s}\left(T^{(1)} - (tu-1)\right) T^{(2)}, \\ M_{+-} &= \frac{T_{+-}}{tu-1}\, T^{(2)} \end{aligned} \tag{9}$$

are free from kinematic singularities where $++$ and $+-$ denote photon helicities same and opposite respectively. The partial wave expansions of these amplitudes can now be expressed as

$$M_{++} = 2\pi \sum_{J=0}^{\infty} (2J+1)\, f_+^J(s)\, d_{00}^J(\theta), \tag{10}$$

$$M_{+-} = 2\pi \sum_{J=2}^{\infty} (2J+1)\, f_-^J(s)\, \frac{d_{20}^J}{1-x^2}(\theta).$$

Since both $\pi^+\pi^-$ and $\pi^0\pi^0$ final states have to be considered, the partial wave amplitudes can be decomposed in the isospin space in the following manner:

$$\begin{aligned} f_\pm^{J,c}(s) &= (+1/\sqrt{3})\, f_\pm^{J,I=0}(s) + (1/\sqrt{6})\, f_\pm^{J,I=2}(s), \\ f_\pm^{J,n}(s) &= (-1/\sqrt{3})\, f_\pm^{J,I=0}(s) + (\sqrt{2/3})\, f_\pm^{J,I=2}(s), \end{aligned} \tag{11}$$

where c and n refer to charged and neutral pions, respectively. Again, the differential cross-section is given by

$$\frac{d\sigma}{d\Omega} = \frac{1}{64\pi^2 s} \sqrt{\frac{s-4}{s}}\, |T|^2. \tag{12}$$

In the range of energies considered here, i.e., $\sqrt{s} < 1{\cdot}4$ GeV we assume that only s and d waves are significant. Thus the explicit forms of the total cross-sections are given by

$$\sigma^c_+ = \alpha \{ | f^{0,0}_+(s) + (1/\sqrt{2}) f^{0,2}_+(s) |^2 + 5 | f^{2,0}_+(s) + (1/\sqrt{2}) f^{2,2}_+(s) |^2 \},$$
$$\sigma^n_+ = \alpha \{ | \sqrt{2} f^{0,2}_+(s) - f^{0,0}_+(s) |^2 + 5 | \sqrt{2} f^{2,2}_+(s) - f^{2,0}_+(s) |^2 \},$$
$$\sigma^c_- = \alpha q^4 \{ 5 | f^{2,0}_-(s) + (1/\sqrt{2}) f^{2,2}_-(s) |^2 \},$$
$$\sigma^n_- = \alpha q^4 \{ 5 | \sqrt{2} f^{2,2}_-(s) - f^{2,0}_-(s) |^2 \} \quad (13)$$

in which

$$\alpha = \pi \{ s(s-4) \}^{\frac{1}{2}}/12.$$

The matrix elements for the decays $S \to \gamma\gamma$ and $S \to \pi\pi$, S being a scalar meson with mass m_S and positive charge conjugation, are given by

$$\langle k_1, \varepsilon_1; k_2, \varepsilon_2 | S \rangle = \frac{G_{S\gamma\gamma}}{m_S^2}\, s \left(\varepsilon_1 \cdot \varepsilon_2 - \frac{\varepsilon_1 \cdot k_2\, \varepsilon_2 \cdot k_1}{k_1 \cdot k_2} \right),$$
$$\langle q_1; q_2, I{=}0 | S \rangle = \sqrt{3}\, g_{S\pi\pi}\, m_S. \quad (14)$$

The corresponding decay rates can be written as

$$\Gamma(S \to \gamma\gamma) = g^2_{S\gamma\gamma} / 16\pi\, m_S,$$
$$\Gamma(S \to \pi\pi) = 3 g^2_{S\pi\pi}\, q / 16\pi. \quad (15)$$

From equations (15), assuming Breit-Wigner formula, one can derive the following result for the decay width (Schierholz *et al* 1972)

$$\Gamma(S \to \gamma\gamma) = \tfrac{1}{128} (m_S^2 - 4)^{\frac{1}{2}}\, m_S^3\, \Gamma_S \times (\mathrm{Im}\, f^{0,0}_+(m_S^2))^2. \quad (16)$$

In the above Γ_S stands for the total width.

3. Dispersion relations and their solutions

The partial wave dispersion relations for the interaction under consideration can be explicitly written down as follows:

$$f^{J,I}_\pm(s) = f^{J,I}_{B\pm}(s) + \frac{1}{\pi} \int_L \frac{\mathrm{Im}\, f^{J,I}_\pm(s')}{s'-s}\, ds' + \frac{1}{\pi} \int_R \frac{\mathrm{Im}\, f^{J,I}_\pm(s')}{s'-s}\, ds'. \quad (17)$$

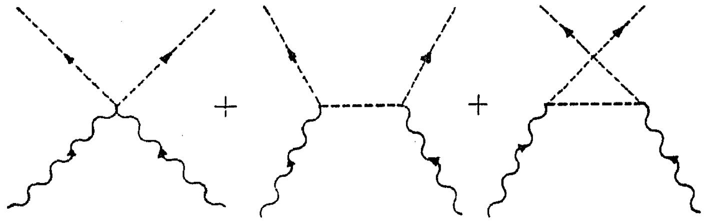

Figure 2. Single pion exchange graphs.

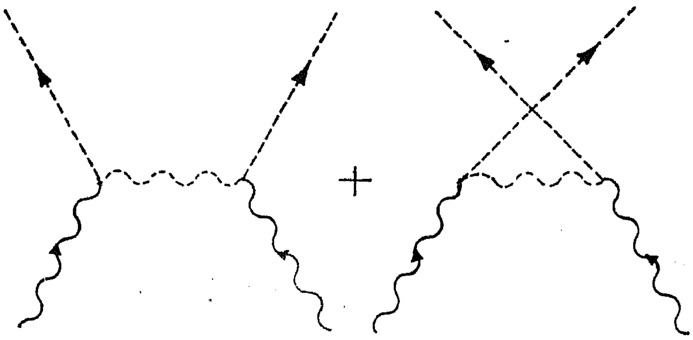

Figure 3. Crossed channel resonance contribution.

Here $f_{B\pm}^{J,\,I}$ are the projections of Born terms for charged pion production (figure 2). In the Born approximation (no strong interactions) the amplitude is $e^2\, N_B^{\mu\nu}\, \varepsilon_{1\mu}\, \varepsilon_{2\nu}$, where

$$N_B^{\mu\nu} = -\,2g^{\mu\nu} + \frac{(2q_1^\mu - k_1^\mu)\,(2q_2^\nu - k_2^\nu)}{2q_1 \cdot k_1} + \frac{(2q_1^\nu - k_2^\nu)\,(2q_2^\mu - k_1^\mu)}{2q_2 \cdot k_1}. \tag{18}$$

After a brief calculation, the Born term contributions come out as

$$T_{++} = -2e^2 s \,/\, (s^2/4 - sq^2\, x^2)$$
$$T_{+-} = -2e^2\, sq^2\, (1-x^2) \,/\, (s^2/4 - sq^2\, x^2). \tag{19}$$

The left-hand cut is approximated assuming that the dominant contributions are from the ρ and ω-vector meson exchanges (figure 3). This gives (Schierholz *et al* 1972, Sundermeyer 1974)

$$T_V = g^2_{V\pi\gamma}\, \epsilon_{\mu\nu\tau\sigma}\, k_1^\mu\, \varepsilon_1^\nu\, \Delta^\tau \frac{g^{\sigma\lambda} - \Delta^\sigma\, \Delta^\lambda / m_V^2}{t - m_V^2}\, \varepsilon_{\lambda\alpha\beta\gamma}\, \Delta^\alpha\, \varepsilon_2^\beta\, k_2^\gamma + (q_1 \longleftrightarrow q_2),$$

where $\Delta = q_1 - k_1$, so that

$$T_V^{(1)} = (g^2_{V\pi\gamma}/4) \left\{ \frac{t(s-u)+1}{t-m_V^2} + \frac{u(s-t)+1}{u-m_V^2} \right\}, \tag{20}$$

$$T_V^{(2)} = -(g^2_{V\pi\gamma}/4) \left\{ \frac{1}{t-m_V^2} + \frac{1}{u-m_V^2} \right\}. \tag{21}$$

We take the vector mesons ρ and ω to be on the α_ρ and α_ω-trajectories, respectively. The left-hand cut integral is expressed by the partial wave projections (10) of the

ρ and ω exchange contributions (21) including a form factor F which has the form (Schierholz *et al* 1972)

$$F = \exp\{(\alpha_v(t) - 1)\,\xi(s)\}. \tag{22}$$

The form of the ξ —function is

$$\xi(s) = \cosh^{-1}(1+(s-4)(9s-4)/32s). \tag{23}$$

Here it is to be remembered that ρ contributes to both charged and neutral pion production while ω contributes only to the neutral pion production. Also in the $\rho\pi\gamma$ vertex the photon must couple through the isoscalar current, whereas in the vertex $\omega\pi\gamma$ it must couple through the isovector current. The amplitudes corresponding to the vector-meson exchanges are denoted by $f^{J,I}_{V\pm}$.

Now let us consider the right-hand cut. Invoking the basic conservation laws, we find that the partial waves with odd angular momenta are absent. Also recent $\pi\pi$-phase shift analyses by the Berkeley group (Protopopescu *et al* 1973) and the CERN—Munich group (Hyams *et al* 1973 and 1975, Grayer *et al* 1974) show that the $\pi\pi$ phase shifts are influenced very much by the opening of inelastic channels around 1 GeV. Therefore, we assume inelastic unitarity for the right-hand cut so that we may write

$$\mathrm{Im} f^{J,I}_{\pm}(s) = [(1-\eta^{I}_{J}(s)\exp(-2i\,\delta^{I}_{J}(s))]\, f^{J,I}_{\pm}(s)/2i. \tag{24}$$

Here δ^{I}_{J} and η^{I}_{J} are the $\pi\pi$ phase shifts and inelasticities, respectively. Then the dispersion relations take the form

$$f^{J,I}_{\pm}(s) = C^{J,I}_{\pm}(s) + \frac{1}{2i\pi}\int_{4}^{\infty} \frac{[1-\eta^{I}_{J}(s')\exp(-2i\,\delta^{I}_{J}(s'))]\; f^{J,I}_{\pm}(s)}{s'-s}\, ds', \tag{25}$$

where

$$C^{J,I}_{\pm}(s) = f^{J,I}_{B\pm}(s) + f^{J,I}_{V\pm}(s). \tag{26}$$

Taking into account the asympotic behaviour of $f^{J,I}_{V\pm}$, it is assumed that the dispersion relation (25) holds without subtraction*. It is also worth pointing out here that Frye and Warnock (1963) have argued on the basis of the dominance of single Regge trajectory at infinite energy that the limits $\eta \to 1$† and $\delta \to n\pi$ are strongly suggestive by the Regge pole hypothesis, even though one has not quite proved that they follow from it. Equation (25) can be solved by the well-known Omnès-Muskhelishvili method (Omnès 1958, Muskhelishvili 1953, Jackson 1961). The solution can be written as

$$f^{J,I}_{\pm}(s) = C^{J,I}_{\pm}(s) + \frac{\exp(u^{I}_{J}(s))}{2i\pi} \times \int_{4}^{\infty} \frac{C^{J,I}_{\pm}(s')\,[\exp(2i\,\delta^{I}_{J}(s')) - \eta^{I}_{J}(s')]\exp(-u^{I}_{J}(s'))}{\eta^{I}_{J}(s')\;(s'-s)}\, ds', \tag{27}$$

*See for example Nielsen (1971) for a discussion of the advantages of an unsubtracted dispersion relation.

†$\eta=0$ corresponds to the so called black-absorber model (Atkinson 1966).

in which

$$u_J^I(s') = \frac{s - s_0}{\pi} \int_4^{\infty} \frac{\delta_J^I(s')}{(s' - s)(s' - s_0)} ds'$$

$$- \frac{(s - s_0)}{\pi} \int_4^{\infty} \frac{\ln \eta_J^I(s')}{(s' - s)(s' - s_0)} ds'. \tag{28}$$

The solution, as determined above, is uncertain up to an arbitrary polynomial P (the presence of which is a manifestation of the well-known *CDD* ambiguity (Castillejo *et al* 1964) whose degree should not be higher than that allowed by the existence of integral in eq. (25). For our case we may assume that this P is identically zero. In other words, we make our solution unique by taking the solution which falls off rapidly as $s \to \infty$. This choice is justified since it has been shown that the choice $P = 0$ is capable of producing a reasonable result as in the case of pion nucleon scattering and pion photoproduction (Adler 1968, Schwela *et al* 1967).

4. Interaction constants

The coupling constants $g_{\omega\pi\gamma}$ and $g_{\rho\pi\gamma}$ are calculated from the experimental values of the decay widths (Particle Data Group 1974) *viz.*,

$$\Gamma(\omega \to \pi\gamma) \approx 0{\cdot}9 \text{ MeV and } \Gamma(\rho \to \pi\gamma) \approx 0{\cdot}4 \text{ MeV}.$$

Correspondingly we have

$$g_{\omega\pi\gamma} \approx 0{\cdot}11 \text{ and } g_{\rho\pi\gamma} \approx 0{\cdot}043.$$

These values approximately satisfy the SU (3) relation

$$g_{\rho\pi\gamma} = (1/3)\, g_{\omega\pi\gamma}.$$

For the Regge trajectories in respect of the ρ and ω mesons we assume the form

$$\alpha_\rho(t) = \alpha_\omega(t) = 0{\cdot}57 + 0{\cdot}96\, t,$$

where the ρ trajectory has been taken from a widely accepted fit to the pion nucleon charge exchange reactions (Höhler *et al* 1966). We incorporate the experimental values of the isospin zero s and d wave phase shifts from the data of the Berkeley group which finally resolves the 'up down' ambiguity in favour of the 'down' solution. Note that Schierholz *et al* have on the other hand made use of a parametrization based on the Breit-Wigner formula. The values of δ_0^2 are taken from Colton *et al* (1971). Also δ_2^2 is taken to be zero since all experimental analyses (Männer 1974) suggest that it is extremely small, at the energies considered here.

5. Discussion of results and conclusions

With the set of parameters given in the previous section we now present our results (figures 4-8). At low energies, the results of Lyth and others (Lyth 1971, 1972, Isaev *et al* 1972) that the Born terms (pion exchange) give the dominant contribution are confirmed.

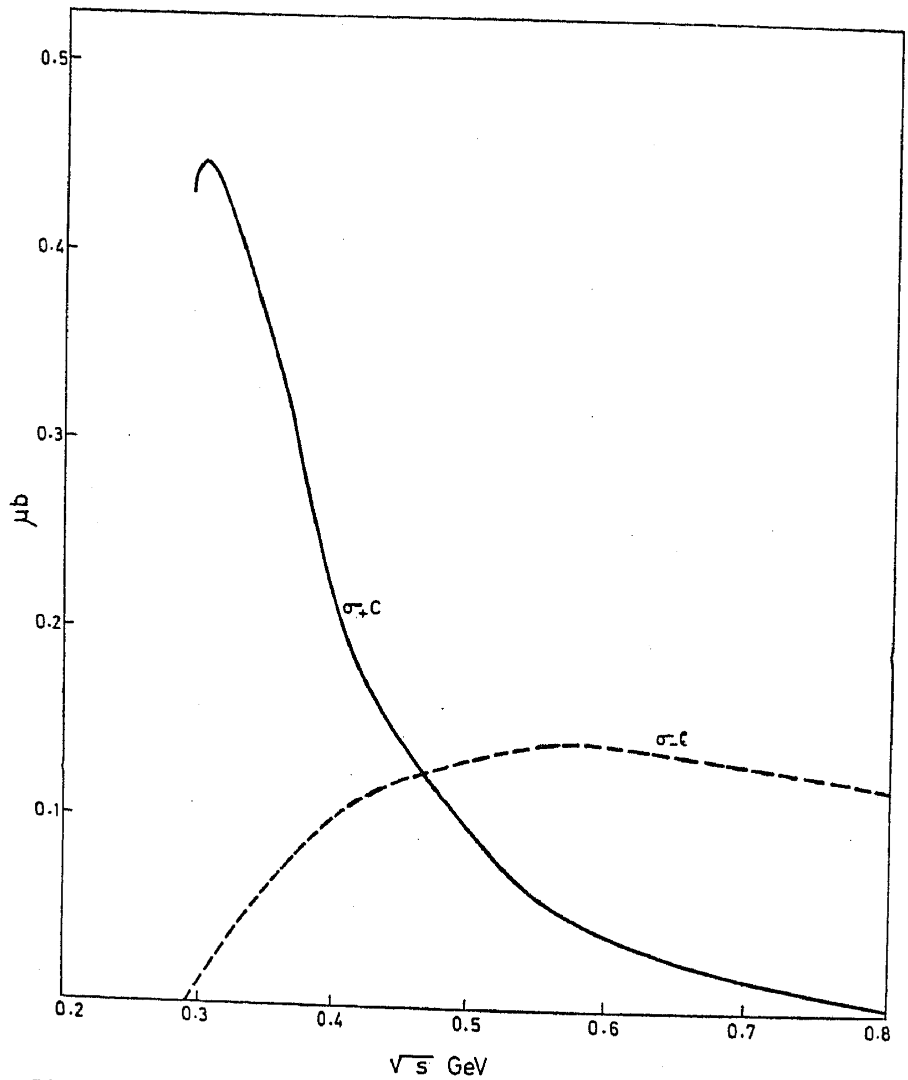

Figure 4. Cross sections with pion Born terms only.

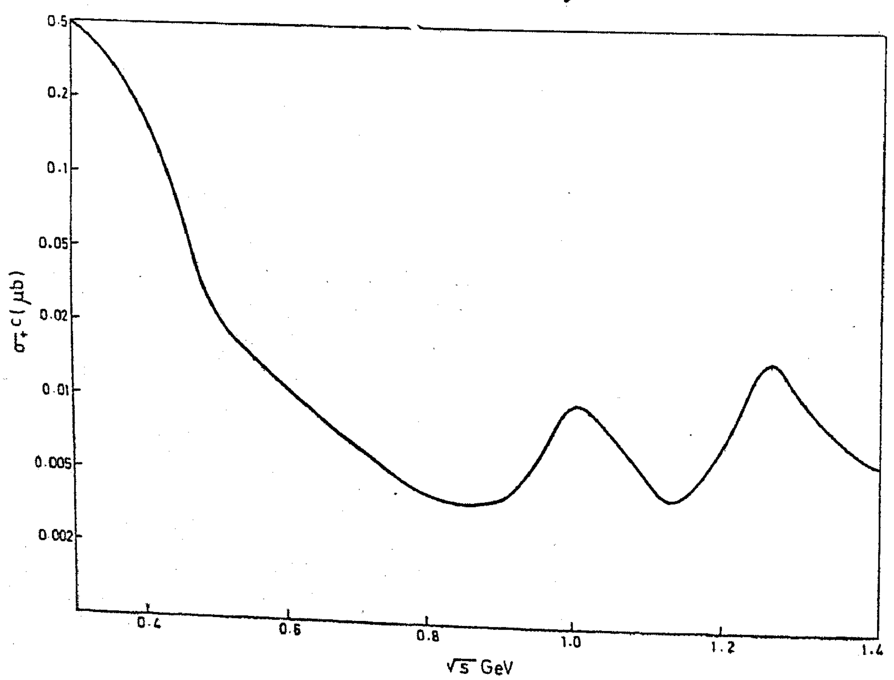

Figure 5. Total cross sections for the production of charged pions when photon helicities are same.

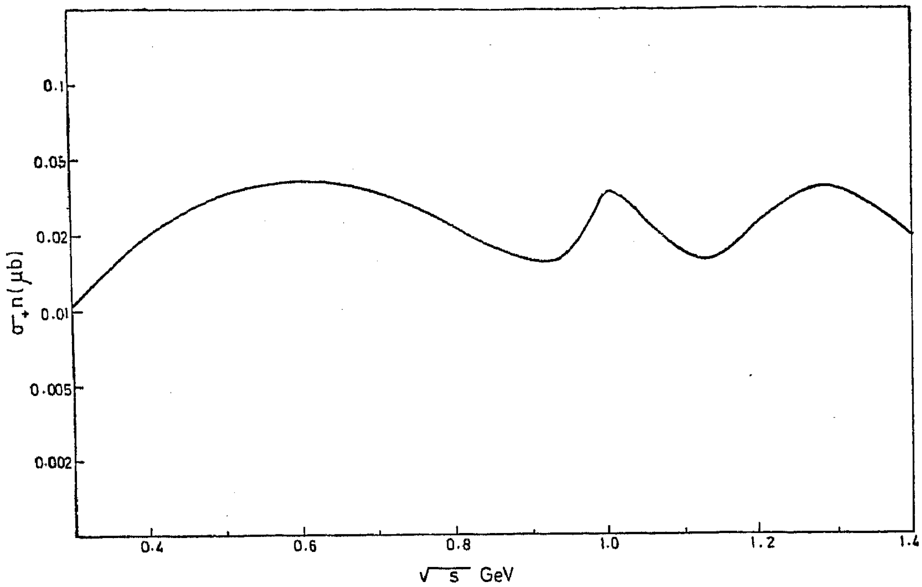

Figure 6. Total cross sections for the production of neutral pions when photon helicities are same.

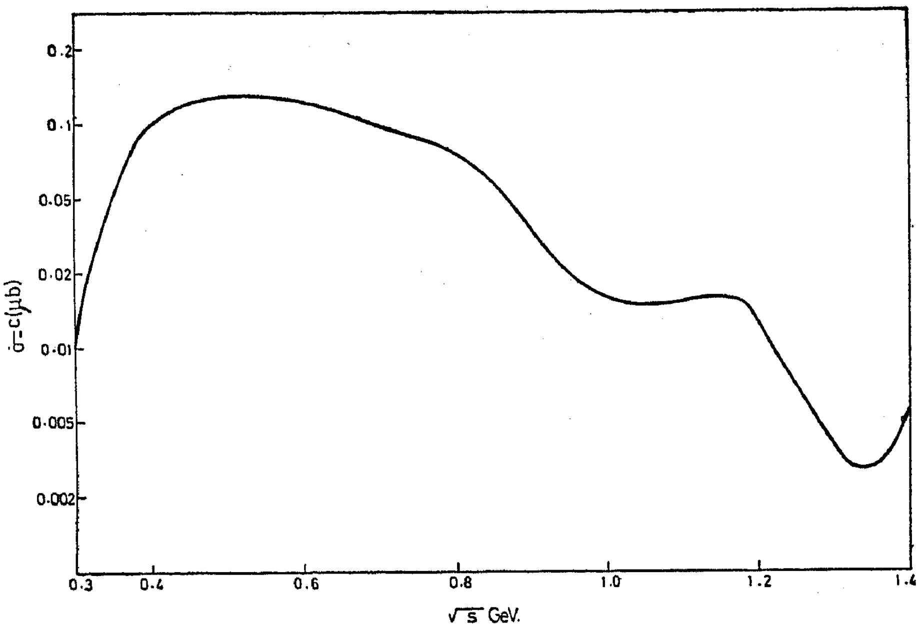

Figure 7. Total cross sections for the production of charged pions when photon helicities are opposite.

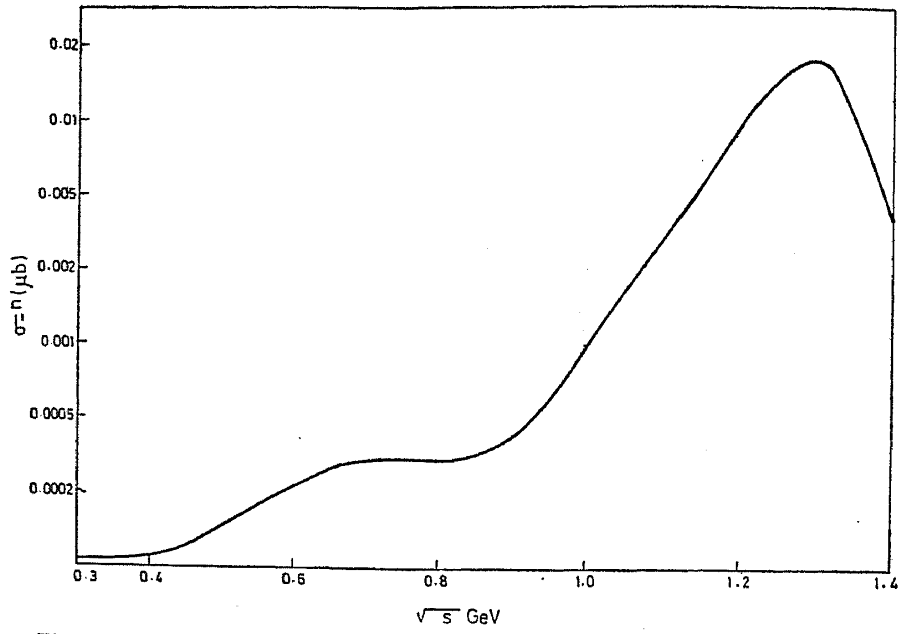

Figure 8. Total cross sections for the production of neutral pions when photon helicities are opposite.

We observe from figures 5 and 6 that both σ_{+}^{c} and σ_{+}^{n} show clear indication of the production of S^* (997) and $f(1270)$ mesons. The structure in figures 7 and 8 is different, which may be accounted for by the fact that contributing partial wave has the minimum angular momentum two, in the case when photon helicities are in the opposite direction (vide eqs. (10) and (13)). These bumps evidently have come from the s and d—partial waves of the $\pi\pi$ system. The bumb at S^*-position is absent in the work of Schierholz *et al* (1972) possibly because of some other parametrization of the s-wave phase shift. Hence we conclude that the reaction

$$\gamma\gamma \to \pi\pi$$

is dominated by the final state $\pi\pi$ interaction. In view of the fact that the existence of S^* meson is often doubted, our results seem to be valuable and useful.
Writing now,

$$\frac{d\sigma}{d\Omega} = \frac{d\sigma_+}{d\Omega} + \frac{d\sigma_-}{d\Omega}$$

we find

$$\frac{d\sigma}{d\Omega} \approx 0{\cdot}038\ \mu b\ at\ \theta = 90^\circ \text{ and } \sqrt{s} = 620 \text{ MeV},$$

to be compared with the only experimental value (Orito *et al* 1974) $d\sigma/d\Omega \approx 0{\cdot}06\ \mu b$. The disagreement in the present numerical results could be due to the large uncertainty in the experimental results or due to our neglecting the contribution of other particles in the crossed channel which can couple to the $\gamma\pi$ system. Another important point to be remembered is that the main peculiarities of the behaviour of the cross section also depend considerably on the particular kind of parametrization used for the scattering phases (Isaev *et al* 1973).

Now if we expect that S^*-meson can have an appreciable decay mode to $\pi\pi$ (which is to be confirmed because of its strong coupling to the $K\bar{K}$ system), then we can calculate its decay width into two photons using eq. (16). In this way, taking $\Gamma(S^* \to \pi\pi) \approx 40$ MeV, and $m_{S^*} \approx 997$ MeV we obtain

$$\Gamma(S^* \to \gamma\gamma) \approx 1 \text{ keV}.$$

The justification for S^* to have an appreciable decay mode to $\pi\pi$ lies in the fact that the phase space available for the S^* meson for its $K\bar{K}$ decay is about one tenth of the phase space available for the $\pi\pi$-channel. Pressuming S^*-meson as the 'mostly octet' isoscalar meson and $\varepsilon(700)$ as the 'mostly singlet' isoscalar state of a possible scalar nonet, Conforto (1974) has shown that a choice of $\varepsilon - S^*$ mixing angle of the order of 37° can give a branching ratio

$$\frac{\Gamma(S^* \to \pi\pi)}{\Gamma(S^* \to \text{all})} = 0{\cdot}75.$$

Combining the anomalous Ward identity for $\theta^\lambda_\lambda - J^\mu - J^\nu$, where θ^λ_λ is the trace of the stress energy momentum tensor and J^μ is the electromagnetic current, with partially conserved dilation current (PCDC) via two isoscalar mesons, Eliezer (1975) has recently estimated the $S^* \to \gamma\gamma$ decay width. In this way he is led to the result:

$$\Gamma(S^* \to \gamma\gamma) \approx 0{\cdot}2 \text{ keV}$$

assuming $\Gamma(S^* \to \pi\pi) \approx 50$ MeV and $m_{S^*} = 980$ MeV, comparable to our result, obtained above.

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Note added in proof

An alternate way to compute the $S^* \to \gamma\gamma$ decay width is as follows:

Using the $q\bar{q}$ bound state description for P waves (Barbieri *et al* 1976 *Phys. Lett.* **B60** 183) we can write

$$\Gamma(O^{++} \to \gamma\gamma) = (9/2)\,(96\alpha^2 e_q^4 / m_S^4)\,|\,\varphi'_P(0)\,|^2\,,$$

where $\alpha = 1/137$, e_q is the charge of the quark q and $|\,\varphi'_P(0)\,|^2 = 0{\cdot}01$ GeV^5 for the ordinary quark bound states. If S^* is made of $\lambda\bar{\lambda}$ ($e_q = -1/3$) and taking $m_{S^*} = 0{\cdot}997$ GeV, we get

$$\Gamma(S^* \to \gamma\gamma) \approx 2{\cdot}65 \text{ keV}.$$

The value of $|\,\varphi'_P(0)\,|$ (see Barbieri *et al* above) is of course much dependant on the choice of the $q\bar{q}$ potential.

One of the present authors (TC) is grateful to F M Renard (Montpellier, France) for pointing out this.